GRAVITATIONAL WAVES FROM PHASE TRANSITIONS OF ACCRETING NEUTRON STARS

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ABSTRACT

We propose that when neutron stars in low-mass X-ray binaries accrete sufficient mass and become millisecond pulsars, the interiors of these stars may undergo phase transitions, which excite stellar radial oscillations. We show that the radial oscillations will be mainly damped by gravitational-wave radiation instead of internal viscosity. The gravitational waves can be detected by the advanced Laser Interferometer Gravitational-Wave Observatory at a rate of about three events per year.

Subject headings: dense matter — gravitation — stars : neutron

1. INTRODUCTION

Gravitational wave astronomy may soon become an observational science, since three gravitational wave experiments, including the Laser Interferometer Gravitational-Wave Observatory (LIGO) (Abramovici et al. 1992), are under construct. In astrophysics, neutron stars are widely believed to be the most promising source of gravitational radiation (for detailed reviews see Thorne 1987, 1995), which from collapse of the cores of massive stars may provide a signature for the features of supernova explosion (Burrows & Hayes 1996), which from starquakes of pulsars may reveal the physics of the stellar interiors (Zimmermann & Szedenits 1979), and which from mergers of binary neutron stars may give information about the equation of state (EOS) of nuclear matter at high densities (Shibata, Nakamura & Oohara 1992, 1993; Rasio & Shapiro 1994; Zhuge, Centrella & McMillan 1994; Davies et al. 1994; Ruffert, Janka & Schäfer 1996). In this Letter we propose a new possible origin of gravitational-wave bursts. We argue that when neutron stars in low-mass X-ray binaries accrete sufficient mass and become millisecond pulsars the interiors of these stars can undergo phase transitions, which excite stellar radial oscillations, producing strong gravitational wave bursts.

2. THE MODEL

2.1. Evolution of Neutron Stars in Low-Mass X-Ray Binaries

According to the standard scenario of evolution of low-mass X-ray binaries (Bhattacharya & van den Heuvel 1991), mass is transferred from the companion to the neutron star, which is spun up to a millisecond period. The mass-transfer from the companion keeps driving the two stars apart while processes such as orbital gravitational radiation or magnetic braking keeps driving the two stars closer. These processes keep the system in a steady mass-transfer state throughout the evolutionary timescale of the companion while the accretion rate of the neutron star is near the Eddington value. Thus, the neutron star can accrete mass $\geq 0.5 M_\odot$ in $\sim 10^8$ years and become a millisecond pulsar (van den Heuvel & Bitzaraki 1995a, b). If we make an assumption that the masses of neutron stars before accretion are $1.4 M_\odot$, which is supported by the current theories of Type II supernova explosion and observations of masses of pulsars (e.g., the Hulse-Taylor binary

system), then the stars in an evolutionary timescale $\geq 10^8$ yrs must become rather massive ones ($\geq 1.8M_\odot$). Now we ask a question: what could possibly occur in the interiors of these massive neutron stars?

To study this question, we first analyze possible EOSs for neutron stars. So far there have been many approaches to determine an EOS for dense matter through the many-body theory of interacting hadrons. Unfortunately, these approaches have given EOSs with different stiffnesses and in turn very different structures of neutron stars. However, the EOSs should be constrained by the observations as follows. First, Link, Epstein & van Riper (1992) used a model-independent approach to analyze the postglitch recovery in four isolated pulsars (Crab, Vela, PSR 0355+54 and PSR 0525+21) which are likely to be isolated $1.4M_\odot$ neutron stars, and concluded that soft EOSs at high densities are ruled out. More detailed analyses of the postglitch curves of the Crab and Vela pulsars also draw similar conclusions (Alpar et al. 1993, 1994). Second, if the EOSs in cores of neutron stars with mass $\sim 1.4M_\odot$ were soft, the massive compact objects after the accretion phase of low-mass X-ray binaries could be black holes (Brown 1988). In fact, these objects have been identified as millisecond pulsars. This means that soft EOSs are unlikely to occur in neutron stars with mass $\sim 1.4M_\odot$. For these two reasons, we can assume that the EOS in neutron stars with mass $\sim 1.4M_\odot$ is moderately stiff to stiff.

The above assumption is consistent with recent theoretical studies of the EOS for dense matter at high densities. First, because of the strong repulsion between nucleons and nucleon holes in the spin-isospin interaction, pion condensation is unlikely in neutron stars (Brown et al. 1988; Baym 1991). Second, the possibility of kaon condensation in dense matter was suggested by Kaplan and Nelson (1986), whose basic idea is that the energy of a negative kaon is lowered by interaction with nucleons. In neutron star matter in beta equilibrium one expects negative kaons to be present if the energy to create one kaon in the matter is less than the electron chemical potential. However, Pandharipande, Pethick and Thorsson (1995) studied kaon-nucleon and nucleon-nucleon correlations in kaon condensation in dense matter and found the kaon energy is much larger than that of the previous studies (see, e.g., Brown et al. 1992; Brown et al. 1994; Thorsson, Prakash & Lattimer 1994; Lee et al. 1995; Thorsson & Wirzba 1995). On the other hand, if we adopt electron chemical potentials of the calculations of the simple parametrized models of Prakash, Ainsworth and Lattimer (1988) which represent a number of more realistic models and compare these potentials with the kaon energies of the Hartree calculations

(Pandharipande et al. 1995) in Figure 1, then we can see that the density for kaon condensation is $\sim$ (5-7)ρ_0 (where ρ_0 is the nuclear density), which significantly exceeds the central densities of neutron stars with mass $\sim 1.4M_\odot$ and with moderately stiff or stiff EOSs. Several authors (e.g., Ellis, Knorren & Prakash 1995; Schaffner & Mishustin 1996; Dai & Cheng 1996) have used the relativistic mean-field approach to study kaon energies in neutron star matter, and have also drawn similar conclusions. Third, it is thought that the density for deconfinement of nuclear matter to two-flavor quark matter is near (6-9)ρ_0 (Baym 1991). Therefore, the recent theoretical studies of dense matter also indicate that pion (or kaon) condensation or quark matter is unlikely to occur in neutron stars with mass $\sim 1.4M_\odot$.

We now turn to the study of what happens in neutron stars with moderately stiff to stiff EOSs when they accrete mass $\geq 0.5M_\odot$, viz., when their masses increase from $1.4M_\odot$ to more than $1.8M_\odot$. It is clear that the central densities of these massive stars are $\sim$(5-7)ρ_0. Once this condition is reached, several physical processes will take place in the stars. When $e^- \rightarrow K^- + \nu_e$, kaon condensation occurs spontaneously in the stellar interior. The appearance of this new phase destroys the hydrostatic equilibrium of the star due to the softening of the EOS of the core matter. This implies a structural transition of the whole star into a new, stable configuration with a condensed core of the new phase. The structural transition — neutron star corequake — occurs on the dynamic timescale of milliseconds, in which the interior temperature may increase to $\sim 10\,\mathrm{MeV}$ both due to the rapid compression of dense matter (Haensel, Denissov & Popov 1990) and due to the energy released by reaction $e^- \rightarrow K^- + \nu_e$. Alternatively, if the central nuclear matter is deconfined into two-flavor quark matter, the quark matter will convert to three-flavor quark matter because strange matter is thought to be more stable than nuclear matter. Thus a strange matter seed is formed in the interior, and subsequently the strange matter will begin to swallow the neutron matter in the surroundings. This process should proceed in a timescale $\sim$ tens of milliseconds due to a detonation mode, and the interior temperature increases to $\geq 10\,\mathrm{MeV}$ because the chemical energy of the two-flavor quark matter is dissipated into thermal energy (Dai, Peng & Lu 1995). The conversion of neutron stars to strange stars has been suggested as a possible origin of cosmological γ-ray bursts (Cheng & Dai 1996). Furthermore, the phase transitions of massive neutron stars to either stars with kaon condensation cores or strange stars can stimulate stellar radial oscillations. In next subsection we will show that these newly born

compact objects may be a strong source of gravitational-wave radiation if their rotation periods are of the order of milliseconds.

2.2. Damping of Radial Oscillations

Radial oscillations are damped not only due to dissipation of the vibration energy of stellar matter into heat but also due to conversion of this energy into gravitational-wave radiation. We first focus on the case in which a massive neutron star undergoes a phase transition to a star containing a kaon condensation core. After $e^- \to K^- + \nu_e$, electron neutrinos are trapped in the interior and form an ideal Fermi-Dirac gas with $\mu_\nu \gg kT$ (where T is the interior temperature) because the neutrino mean free path is much less than the stellar radius for $KT \sim 10\,\mathrm{MeV}$. Subsequent evolution of the newborn star, which is analogous to that of a protoneutron star formed from supernova explosion (Burrows & Lattimer 1986), can be divided into three stages. We neglect gravitational radiation. The first stage is deleptonization, whose timescale (τ_{d1}) is of the order of 1 s (Sawyer & Soni 1979), in which the interior temperature may increase significantly. In this stage the radial oscillations are damped through the following reaction

$$\tilde{p} + e^- \leftrightarrow \tilde{n} + \nu_e \,, \tag{1}$$

where $\tilde{p}$ and $\tilde{n}$ represent quasiprotons and quasineutrons respectively. Using the reaction matrix element of Brown et al. (1988), we have derived the net rate per baryon for this reaction at nonequilibrium,

$$\Gamma_1 = \frac{1}{24\pi^5\hbar^{10}c^5} G_F^2 \cos^2\theta_C \cos^2(\theta/2)(1+3g_A^2) m_n^* m_p^* \mu_e \mu_\nu^2 \Delta\mu [\Delta\mu^2 + (2\pi kT)^2] \,, \tag{2}$$

where G_F is the weak coupling constant, θ_C is the Cabibbo angle, θ is the chiral angle for kaon condensation ($\theta \leq 60^{\mathrm{o}}$) (Brown et al. 1994), g_A is the Gamow-Teller coupling constant, m_n^* and m_p^* are the effective nucleon masses, and $\Delta\mu = \mu_p + \mu_e - \mu_n - \mu_\nu$ with μ_i being the chemical potential of particle i. In deriving equation (2), we have assumed $kT/c \ll |\mathbf{P}_\nu| \ll |\mathbf{P}_n|$ with $\mathbf{P}_\nu$ and $\mathbf{P}_n$ being the neutrino and neutron momenta respectively. As the steps shown by Dai & Lu (1996), using this reaction rate, we have further derived the bulk viscosity as

$$\eta_1 \simeq 1.8 \times 10^{25} \cos^{-2}(\theta/2) Y_e^{-1/3} Y_\nu^{2/3} Y_n^{4/3} \left(\frac{\rho}{\rho_0}\right)^{5/3} \left(\frac{kT}{1\mathrm{MeV}}\right)^{-2} \mathrm{g\,cm^{-1}\,s^{-1}} \,, \tag{3}$$

where Y_e, Y_ν and Y_n are the particle concentrations, and ρ is the stellar density. Therefore, the damping timescale (Sawyer 1980) is given by

$$\tau_{v1} = \frac{1}{30}\rho R^2 \eta_1^{-1} \tag{4}$$

where R is the stellar radius. Inserting equation (3) into this equation, we obtain

$$\tau_{v1} \simeq 0.52\cos^2(\theta/2) Y_e^{1/3} Y_\nu^{-2/3} Y_n^{-4/3} R_6^2 \left(\frac{\rho}{\rho_0}\right)^{-2/3} \left(\frac{kT}{1\mathrm{MeV}}\right)^2 \ \mathrm{s}\,, \tag{5}$$

where R_6 is in units of 10^6 cm. For the typical particle concentrations and $\rho \sim 6\rho_0$, $R \sim 10^6$ cm and $kT \sim 10$ MeV, $\tau_{v1} \sim 12 Y_e^{1/3} Y_\nu^{-2/3} Y_n^{-4/3}\,\mathrm{s} > 12\,\mathrm{s} \gg \tau_{d1}$.

In the second stage, the stellar interior has practically no trapped lepton-number excess, as compared with catalyzed matter. The diffusion of $\nu_e\bar{\nu}_e$ is then driven by the temperature gradient. Locally, the equilibrium distribution function of $\nu_e\bar{\nu}_e$ can be approximated by the Fermi-Dirac one with a zero chemical potential. The neutrino diffusion timescale (τ_{d2}) is of the order of 40 s (Sawyer & Soni 1979). Using the analogy of deriving the bulk viscosity of (Haensel & Zdunik 1992), we obtain the damping timescale

$$\tau_{v2} \simeq 1.8 \times 10^2 \cos^2(\theta/2) Y_e^{-1/3} R_6^2 \left(\frac{\rho}{\rho_0}\right)^{2/3} \ \mathrm{s}\,, \tag{6}$$

Clearly, $\tau_{v2} \gg \tau_{d2}$.

Third, after the neutrino diffusion, the temperature decreases to ~ 1 MeV, and neutrinos can escape freely from the star. The damping timescale becomes

$$\tau_{v3} \simeq 2.6 \times 10^{-5} \cos^{-2}(\theta/2) Y_e^{-1/3} R_6^2 \left(\frac{\rho}{\rho_0}\right)^{2/3} \left(\frac{kT}{1\mathrm{MeV}}\right)^{-4} \left(\frac{\omega}{10^4\mathrm{s}^{-1}}\right)^2 \ \mathrm{s}\,, \tag{7}$$

where ω is the oscillation frequency. Therefore, we can conclude that the radial oscillations are not damped very efficiently by the bulk viscosity until the stage at which neutrinos escape freely.

We now consider gravitational radiation from rapidly spinning and oscillating neutron stars. The timescale for this process (Chau 1967) is

$$\tau_g \simeq 0.41 M_{1.8}^{-1} R_6^{-2} \left(\frac{P}{2\mathrm{ms}}\right)^4 \ \mathrm{s}\,, \tag{8}$$

where $M_{1.8}$ is the stellar mass in units of $1.8M_\odot$, and P is the stellar rotation period. Here we have assumed that nucleons in the stellar interior are nonrelativistic and degenerate, and thus the adiabatic index is equal to 5/3. When $\tau_g \leq \tau_{d1}+\tau_{d2}$, viz.,

$$P \leq 6.2\,\bar{\tau}_d^{1/4} M_{1.8}^{1/4} R_6^{1/2}\ \mathrm{ms}\,, \tag{9}$$

where $\bar{\tau}_d = (\tau_{d1}+\tau_{d2})/40\mathrm{s}$, the gravitational radiation can damp the radial oscillations very efficiently. The frequency of the gravitational waves is euqal to $\omega = 2\pi/\tau$, where τ is close to $5\times10^{-4}\,$s for a typical neutron-star model (Glass & Lindblom 1983) (corrected for the gravitational redshift) with mass $M \sim 1.8M_\odot$ and radius $R \sim 10^6\,$cm. The strength of the waves can be estimated using the quadrupole approximation to the Einstein field equations (Thorne 1987). This approximation shows that the gravitational strain is given by

$$h \simeq \frac{G}{c^4}\frac{\ddot{Q}}{r} \sim 1.5\times10^{-23} R_6^5 \left(\frac{\rho}{6\rho_0}\right)\left(\frac{\alpha}{0.1}\right)\left(\frac{P}{2\mathrm{ms}}\right)^{-2}\left(\frac{r}{100\mathrm{Mpc}}\right)^{-1}\,, \tag{10}$$

where Q is the source's quadrupole moment (Chau 1967), r is the distance of the source from Earth, and α is the relative oscillation amplitude ($=\delta R/R$). Particularly, in the case of phase transitions of neutron stars with stiff EOSs to stars with kaon condensation cores, α is expected to be about 0.1. Furthermore, the characteristic gravitational strain is

$$h_c \simeq h\sqrt{n} \sim 4.3\times10^{-22} M_{1.8}^{-1/2} R_6^4 \left(\frac{\rho}{6\rho_0}\right)\left(\frac{\alpha}{0.1}\right)\left(\frac{\tau}{0.5\mathrm{ms}}\right)^{-1/2}\left(\frac{r}{100\mathrm{Mpc}}\right)^{-1}\,, \tag{11}$$

where n is the number of cycles of gravitational waves in the duration $\sim \tau_g$. The observed number of low-mass X-ray binaries in our Galaxy is $\sim 10^2$ (van Paradijs 1995), and thus the rate for phase transitions of their neutron stars is $\sim 10^{-6}\,\mathrm{yr}^{-1}$ because the typical accretion timescale is $\sim 10^8$ years. An alternative estimation based on the number of millisecond pulsars and their life time also gives a similar value (Cheng & Dai 1996). Therefore, the rate detected by the advanced LIGO detector (Abramovici et al. 1992) is estimated to be about three events per year.

Alternatively, neutron stars in low-mass X-ray binaries may accrete sufficient mass to convert into strange stars, as suggested by Cheng & Dai (1996). Next we would discuss this case. First, the timescale for damping radial oscillations due to bulk viscosity is $\geq 10\,$s

for a high temperature $\geq 10\,\mathrm{MeV}$ (Madsen 1992; Dai & Lu 1996), so the gravitational radiation damping mechanism is also more efficient in this case. Second, for the conversion of neutron stars with stiff EOSs to strange stars, the relative oscillation amplitude (α) may not be less than 0.1. Third, the density for deconfinement may be larger than that for kaon condensation, and thus the oscillation frequency of strange stars is larger. If gravitational waves discussed in this work are observed by the advanced LIGO in the future, then observations, in principle, can distinguish between these two kinds of phase transitions. In addition, gravitational radiation for the case in which neutron stars convert to strange stars is likely to occur together with cosmological γ-ray bursts, because the fireballs formed during the conversion have very low baryon contamination (Cheng & Dai 1996).

3. CONCLUSIONS

We in this Letter have suggested that during the evolution of neutron stars in low-mass X-ray binaries the stars may undergo phase transitions to stars containing kaon condensation cores or strange stars when these neutron stars accrete sufficient mass from their companions. The phase transitions can excite stellar radial oscillations, which produce strong gravitational wave bursts if the stellar rotation periods are of the order of milliseconds. The study of such gravitational radiation may provide information for the high-density EOS and the physics of phase transitions. In addition, the rate detected by the advanced LIGO is estimated to be about three events per year.

Gravitational waves from colliding neutron stars must have a continuous spectrum, but waves from phase transitions of neutron stars appear to have a delta-function like spectrum. This signature is expected to be confirmed by very near future observations .

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FIGURE CAPTION

FIG. 1. Energy of a single negative kaon in neutron star matter and electron chemical potential as functions of density. The three solid lines are kaon energies from the Hartree calculations of Pandharipande et al. (1995) for square wells of radii $R = 1$ and $0.7\,\mathrm{fm}$, and for a Yukawa potential. The dashed lines labeled by PAL1, PAL2 and PAL3 represent electron chemical potentials calculated from the parametrized models of Prakash et al. (1988) corresponding to three different forms of $F(u)$ which parameterizes the potential contribution to the symmetry energy.